

Electronic control of the spin-wave damping in a magnetic insulator

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It is demonstrated that the decay time of spin-wave modes existing in a magnetic insulator can be reduced or enhanced by injecting an in-plane dc current, I_{dc} , in an adjacent normal metal with strong spin-orbit interaction. The demonstration rests upon the measurement of the ferromagnetic resonance linewidth as a function of I_{dc} in a 5 μm diameter YIG(20nm)|Pt(7nm) disk using a magnetic resonance force microscope (MRFM). Complete compensation of the damping of the fundamental mode is obtained for a current density of $\sim 3 \cdot 10^{11} \text{A.m}^{-2}$, in agreement with theoretical predictions. At this critical threshold the MRFM detects a small change of static magnetization, a behavior consistent with the onset of an auto-oscillation regime.

The spin-orbit interaction (SOI) [1–3] has been recently shown to be an interesting and useful addition in the field of spintronics. This subject capitalizes on adjoining a strong SOI normal metal next to a thin magnetic layer [4]. The SOI converts a charge current, J_c , to a spin current, J_s , with an efficiency parametrized by Θ_{SH} , the spin Hall angle [5, 6]. Recently, it was demonstrated experimentally that the spin current produced in this way can switch the magnetization in a dot [7, 8] or can partially compensate the damping [9–11], allowing the lifetime of propagating spin-waves [12] to be increased beyond their natural decay time, τ . These two effects open potential applications in storage devices and in microwave signal processing.

The effect is based on the fact that the spin current J_s exerts a torque on the magnetization, corresponding to an effective damping $\Gamma_s = \gamma J_s / (t_{FM} M_s)$, where t_{FM} is the thickness of the magnetic layer, M_s its spontaneous magnetization, and γ the gyromagnetic ratio. In the case of metallic ferromagnets [13–15], it was established that Γ_s can fully compensate the natural damping $1/\tau$ at a critical spin current J_s^* , which determines the onset of auto-oscillation of the magnetization:

$$J_s^* = -\frac{1}{\tau} \frac{t_{FM} M_s}{\gamma}. \quad (1)$$

An important benefit of the SOI is that J_c and J_s are linked through a cross-product, allowing a charge current flowing in-plane to produce a spin current flowing out-of-plane. Hence it enables the transfer of spin angular momentum to non-metallic materials and in particular to insulating oxides, which offer improved performance compared to their metallic counterparts. Among all oxides, Yttrium Iron Garnet (YIG) holds a special place for having the lowest known spin-wave (SW) damping factor. In 2010, Kajiwara *et al.* reported on the effi-

cient transmission of spin current through the YIG|Pt interface [16]. It was shown that J_s produced by the excitation of ferromagnetic resonance (FMR) in YIG can cross the YIG|Pt interface and be converted into J_c in Pt through the inverse spin Hall effect (ISHE). This finding was reproduced in numerous experimental works [17–23]. In the same paper, the reciprocal effect was also reported as J_s produced in Pt by the direct spin Hall effect (SHE) could be transferred to the 1.3 μm thick YIG, resulting in damping compensation. However, attempts to directly measure the expected change of the resonance linewidth of YIG as a function of the dc current have so far failed [21, 22] [24]. This is raising fundamental questions about the reciprocity of the spin transparency, T , of the interface between a metal and a magnetic insulator. This coefficient enters in the ratio between J_c in Pt and J_s in YIG through:

$$J_s = T \Theta_{SH} \frac{\hbar}{2e} J_c, \quad (2)$$

where e is the electron charge and \hbar the reduced Planck constant. T depends on the transport characteristics of the normal metal as well as on the spin-mixing conductance $G_{\uparrow\downarrow}$, which parametrizes the scattering of the spin angular momentum at the YIG|Pt interface [25].

At the heart of this debate lies the exact value of the threshold current. The lack of visible effects reported in Refs.[21, 22], although inconsistent with [16], is coherent with the estimation of the threshold current of $10^{11-12} \text{A.m}^{-2}$ using Eqs.(1) and (2) and typical parameters for the materials [26]. This theoretical current density is at least one order of magnitude larger than the maximum J_c that could be injected in the Pt so far. Importantly, the previous reported experiments were performed on large (millimeter sized) structures, where many nearly degenerate SW modes compete for

TABLE I. Transport and magnetic properties of the Pt and bare YIG layers, respectively from Ref.[31] and Ref.[22].

Pt	t_{Pt} (nm)	σ ($\Omega^{-1} \cdot \text{m}^{-1}$)	λ_{SD} (nm)	Θ_{SH}
	7	$5.8 \cdot 10^6$	3.5	0.056
YIG	t_{YIG} (nm)	$4\pi M_s$ (G)	γ ($\text{rad} \cdot \text{s}^{-1} \cdot \text{G}^{-1}$)	α_0
	20	$2.1 \cdot 10^3$	$1.79 \cdot 10^7$	$2.3 \cdot 10^{-4}$

feeding from the same dc source of angular momentum, a phenomenon that could become self-limiting and prevent the onset of auto-oscillations [11]. To isolate a single candidate mode, we have recently reduced the lateral dimensions of the YIG pattern, as quantization results in increased frequency gaps between the dynamical modes [27]. This requires to grow very thin films of high quality YIG [23, 28–30]. Benefiting from our progress in the epitaxial growth of YIG films by pulsed laser deposition (PLD) [22], we propose to study the FMR linewidth as a function of the dc current in a micron-size YIG|Pt disk.

FIG.1 shows a schematic of the experimental setup. A YIG|Pt disk of $5 \mu\text{m}$ in diameter is connected to two Au contact electrodes (see the microscopy image) across which a positive voltage generates a current flow J_c along the $+\hat{x}$ -direction. The microdisk is patterned out of a 20 nm thick epitaxial YIG film with a 7 nm thick Pt layer sputtered on top. The YIG and Pt layers have been fully characterized in previous studies [22, 31]. Their characteristics are reported in Table I.

The sample is mounted inside a room temperature magnetic resonance force microscope (MRFM) which detects the SW absorption spectrum mechanically [32–34]. The excitation is provided by a stripline (not shown in the sketches of FIG.1) generating a linearly polarized microwave field h_1 along the \hat{x} -direction. The detection is based on monitoring the deflection of a mechanical cantilever with a magnetic Fe particle affixed to its tip, coupled dipolarly to the sample. The FMR spectrum is obtained by recording the vibration amplitude of the cantilever while scanning the external bias magnetic field, H_0 , at constant microwave excitation frequency, $f = \omega/(2\pi)$ [35]. The MRFM is placed between the poles of an electromagnet, generating a uniform magnetic field, H_0 , which can be set along \hat{y} or \hat{z} (*i.e.*, perpendicularly to both h_1 and J_c).

We start by measuring the effect of a dc current, I_{dc} , on the FMR spectra when the disk is magnetized in-plane by a magnetic field along the $+\hat{y}$ -direction (positive field). The spectra recorded at $f = 6.33 \text{ GHz}$ are shown in FIG.1a in red tones. The middle row shows the absorption at zero current. The MRFM signal corresponds to a variation of the static magnetization of about 2 G, *i.e.*, a precession cone of 2.5° . As the electrical current is varied, we observe very clearly a change of the linewidth. At negative current, the linewidth decreases,

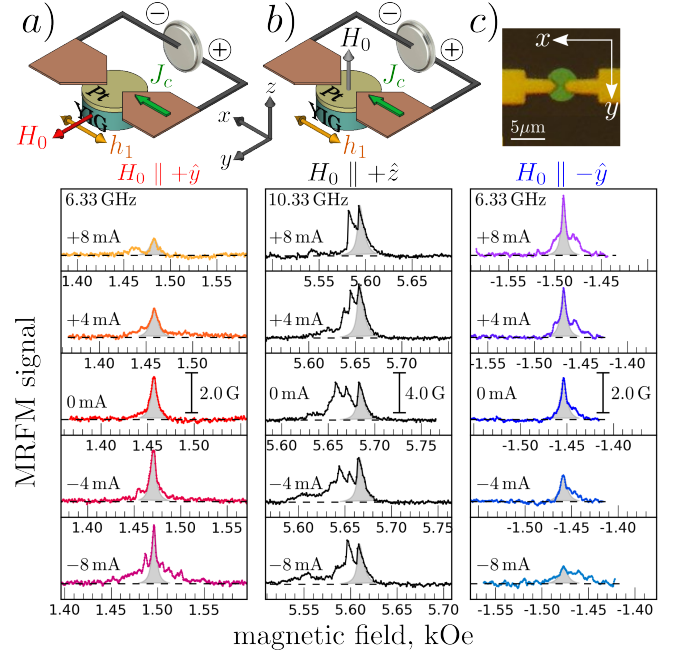


FIG. 1. (Color online) MRFM spectra of the YIG|Pt microdisk as a function of current for different field orientations: a) $H_0 \parallel +\hat{y}$ at $f = 6.33 \text{ GHz}$ (red tone); b) $H_0 \parallel +\hat{z}$ at $f = 10.33 \text{ GHz}$ (black); c) $H_0 \parallel -\hat{y}$ at $f = 6.33 \text{ GHz}$ (blue tone). The highest amplitude mode is used for linewidth analysis (shaded area). Field axes are shifted so as to align the peaks vertically. In-plane and out-of-plane field orientations are sketched above. The top right frame is a microscopy image of the sample.

to reach about half the initial value at $I_{\text{dc}} = -8 \text{ mA}$. This decrease is strong enough so that the individual modes can be resolved spectroscopically within the main peak. Concomitantly the amplitude of the MRFM signal increases. The opposite behavior is observed when the current polarity is reversed. At positive current, the linewidth increases to reach about twice the initial value at $I_{\text{dc}} = +8 \text{ mA}$, and the amplitude of the signal decreases.

$I_{\text{dc}} = \pm 12 \text{ mA}$ is the maximum current that we have injected in our sample to avoid irreversible effects. We estimate from the Pt resistance, the sample temperature to be 90° C at the maximum current. This Joule heating reduces $4\pi M_s$ at a rate of 4.8 G/K , which results in an even shift of the resonance field towards higher field [36].

In FIG.1b, we show the FMR spectra at $f = 10.33 \text{ GHz}$ in the perpendicular geometry, *i.e.*, H_0 is along \hat{z} . In contrast to the previous case, the linewidth does not change with current. This is expected as no net spin transfer torque is exerted by the spin current on the precessing magnetization in this configuration. Note that due to Joule heating, the spectrum now shifts towards lower field due to the decrease of M_s as the current increases.

We now come back to the in-plane geometry, but this time, the magnetic field is reversed compared to FIG.1a,

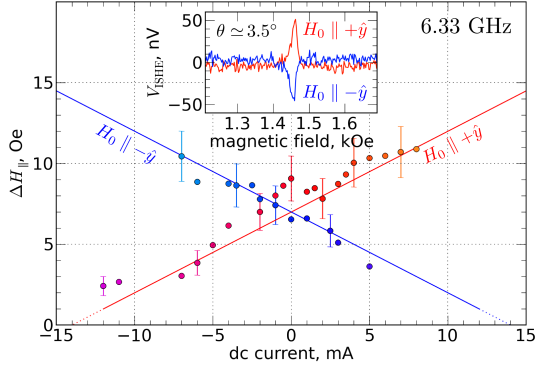


FIG. 2. (Color online) Variation of the full linewidth $\Delta H_{||}$ measured at 6.33 GHz as a function of I_{dc} for $H_0 \parallel +\hat{y}$ (red) and $H_0 \parallel -\hat{y}$ (blue). Inset: detection of V_{SHE} as a function of H_0 at $f = 6.33$ GHz and $I_{dc} = 0$.

i.e., applied along $-\hat{y}$ (negative field). The corresponding spectra are presented in FIG.1c using blue tones. As expected for the symmetry of the SHE, the observed behavior is inverted with respect to FIG.1a: a positive (negative) current now reduces (broadens) the linewidth.

We report in FIG.2 the values of $\Delta H_{||}$, the full linewidth measured in the in-plane geometry, as a function of current. The data points follow approximately a straight line, whose slope ± 0.5 Oe/mA reverses with the direction of H_0 along $\pm \hat{y}$ and whose intercept with the abscissa axis occurs at $I_{6.33 \text{ GHz}}^* = \mp 14$ mA. Moreover, we emphasize that the variation of linewidth covers about a factor five on the full range of current explored.

The inset of FIG.2 shows the inverse spin Hall voltage V_{SHE} measured at $I_{dc} = 0$ mA and $f = 6.33$ GHz. This voltage results from the spin current produced by spin pumping from YIG to Pt and its subsequent conversion into charge current by ISHE [16]. Its sign changes with the direction of the bias magnetic field, as shown by the blue and red V_{SHE} spectra. This observation confirms that a spin current can flow from YIG to Pt and that damping reduction occurs for a current polarity corresponding to a negative product of V_{SHE} and I_{dc} .

To gain more insight into these results, we now analyze the frequency dependence of the full linewidth at half maximum for three values of dc current (0, ± 6 mA) for both the out-of-plane and in-plane geometries. We start with the out-of-plane data, plotted in FIG.3a. The dispersion relation displayed in the inset follows the Kittel law, $\omega = \gamma(H_0 - 4\pi N_{eff}M_s)$, where N_{eff} is an effective demagnetizing factor close to 1 [37, 38]. The linewidth ΔH_{\perp} increases linearly with frequency along a line that intercepts the origin, a signature that the resonance is homogeneously broadened [27]. In this geometry, the Gilbert damping coefficient is simply $\alpha = \gamma \Delta H_{\perp} / (2\omega) = 1.1 \cdot 10^{-3}$ and the relaxation time $\tau = 1/(\alpha\omega)$. We also report on this figure the fact that at 10.33 GHz, $\Delta H_{\perp} = 7$ Oe is independent of the current (see FIG.1b).

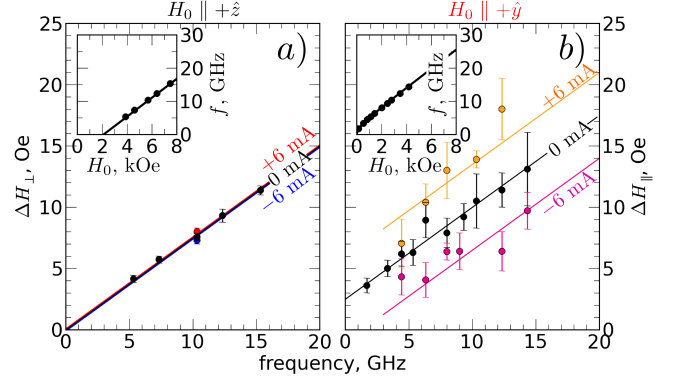


FIG. 3. (Color online) Frequency dependence of the linewidth for three values of the dc current (0, ± 6 mA) a) in the perpendicular geometry and b) in the parallel geometry. Insets show the corresponding dispersion relations $f(H_0)$.

The damping found in our YIG|Pt microdisk is significantly larger than the one measured in the bare YIG film $\alpha_0 = 2.3 \cdot 10^{-4}$ (cf. Table I). This difference is due to the spin pumping effect, and enables to determine the spin-mixing conductance of our YIG|Pt interface through [39, 40]:

$$\alpha = \alpha_0 + \frac{\gamma \hbar}{4\pi M_s t_{YIG}} \frac{G_{\uparrow\downarrow}}{G_0}, \quad (3)$$

where $G_0 = 2e^2/h$ is the quantum of conductance. The measured increase of almost $9 \cdot 10^{-4}$ for the damping corresponds to $G_{\uparrow\downarrow} = 1.5 \times 10^{14} \Omega^{-1} \text{m}^{-2}$, in agreement with a previous determination made on similar YIG|Pt nanodisks [27]. This value allows us to estimate the spin transparency of our interface [25], $T = G_{\uparrow\downarrow} / (G_{\uparrow\downarrow} \coth(t_{Pt}/\lambda_{sd}) + \sigma / (2\lambda_{sd})) \simeq 0.15$, where σ is the Pt conductivity and λ_{sd} its spin-diffusion length. Moreover, the spin-mixing conductance can be used to analyze quantitatively the dc ISHE voltage produced at resonance [21, 41, 42]. Using the parameters of Table I and the value of $G_{\uparrow\downarrow}$, we find that the 50 nV voltage measured in the inset of FIG.2 is produced by an angle of precession $\theta \simeq 3.5^\circ$, which lies in the expected range.

We now turn to the in-plane data, presented in FIG.3b. The dispersion relation plotted in the inset follows the Kittel law $\omega = \gamma \sqrt{H_0(H_0 + 4\pi N_{eff}M_s)}$. In this case, $1/\tau = \alpha(\partial\omega/\partial H_0)(\omega/\gamma)$. For $I_{dc} = 0$ mA the slope of the linewidth vs. frequency is exactly the same as that in the perpendicular direction $\alpha = 1.1 \cdot 10^{-3}$. For this geometry, however, the line does not intercept the origin, indicating a finite amount of inhomogeneous broadening $\Delta H_0 = 2.5$ Oe, *i.e.*, the presence of several modes within the resonance line. Setting I_{dc} to ± 6 mA shifts $\Delta H_{||}$ by ± 3 Oe independently of the frequency, which is consistent with the rate of 0.5 Oe/mA reported at 6.33 GHz in FIG.2. In fact, in the presence of the effective damping

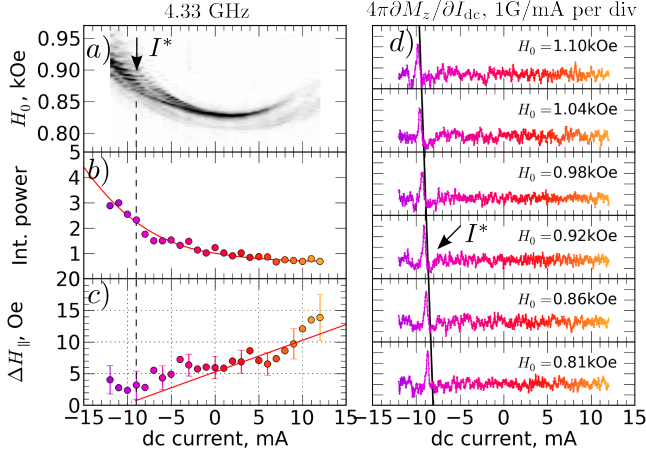


FIG. 4. (Color online) a) Density plot of the MRFM spectra at 4.33 GHz vs. field and current $I_{dc} \in [-12, +12]$ mA. The color scale represents $4\pi\Delta M_z$ (white: 0 G, black: 1.5 G). b) Evolution of integrated power vs. I_{dc} . c) Dependence of linewidth on I_{dc} . d) Differential measurements of M_z (I_{dc} modulated by 0.15 mA_{pp}, no rf excitation) vs. I_{dc} at six different values of the in-plane magnetic field.

Γ_s , the linewidth of the resonance line varies as

$$\Delta H_{\parallel} = \Delta H_0 + 2\alpha \frac{\omega}{\gamma} + 2 \frac{J_s}{M_s t_{YIG}}. \quad (4)$$

This expression is valid when $(\partial\omega/\partial H_0) \simeq \gamma$, *i.e.*, at large enough field or frequency (see inset of FIG.3b). It describes appropriately the experimental data on the whole frequency range measured.

In order to investigate the autonomous dynamics of the YIG layer and exceed the compensation current, I^* , we now perform measurements at lower excitation frequency, where the threshold current is estimated below 12 mA. In FIG.4a, we present a density plot of the MRFM spectra acquired at 4.33 GHz as a function of the in-plane magnetic field and I_{dc} through the Pt. The measured signal is clearly asymmetric in I_{dc} . At positive current, it broadens and its amplitude decreases, almost disappearing above +8 mA, whereas at negative current, it becomes narrower and the amplitude is maximal at $I_{dc} < -10$ mA.

The power integrated over the full field range normalized by its value at 0 mA and the linewidth variation vs. I_{dc} are plotted in FIGs.4b and 4c, respectively. The normalized integrated power varies by a factor of five from +12 mA to -12 mA following an inverse law on I_{dc} (see continuous line), which is consistent with the spin transfer effect [11, 43]. The linewidth varies roughly linearly with I_{dc} : it increases from 6 Oe at 0 mA up to 14 Oe at +12 mA and it reaches a minimum value close to 2 Oe between -8 and -11 mA. It is interesting to note that this happens in a region of the density plot where the evolution of the signal displays some kind of discontinuity, with the appearance of several high amplitude peaks

in the spectrum (see arrow in FIG.4a). We tentatively ascribe this feature to the onset of auto-oscillations in the YIG layer, namely, one or several dynamical modes have their relaxation compensated by the injected spin current and are destabilized [16].

To confirm this hypothesis, we present in FIG.4d results of an experiment where no rf excitation is applied to the system. Here, the dc current is modulated at the MRFM cantilever frequency by $\delta I = 0.15$ mA_{pp} and the induced δM_z is probed as a function of I_{dc} . This experiment thus provides a differential measurement $\partial M_z/\partial I_{dc}$ of the magnetization (in analogy with dV/dI measurements in transport experiments). At $H_0 = 0.92$ kOe, a peak in $\partial M_z/\partial I_{dc}$ is measured around -9 mA. It corresponds to a variation of $4\pi\delta M_z \simeq 0.5$ G, *i.e.*, a change of the angle of precession by 1.3° induced by the modulation of current. Moreover, this narrow peak observed in $\partial M_z/\partial I_{dc}$ shifts linearly in dc current with the applied magnetic field, from -8 mA at 0.81 kOe to -10 mA at 1.1 kOe (see the continuous straight line in FIG.4d), in agreement with the expected behavior of the threshold current Eq.(1).

Hence, FIG.4 presents a set of data consistent with the determination of a critical current of $I^* = -9$ mA at $H_0 = 0.92$ kOe, corresponding to $J_c^* \simeq 3 \cdot 10^{11}$ A.m⁻², in agreement with the value of $2 \cdot 10^{11}$ A.m⁻² expected from Eqs.(1) and (2) and the parameters of our system. Nevertheless, the destabilization of dynamical modes is rather small, as the jump of resonance field at I^* (due to reduction of the magnetization) does not exceed the linewidth. We suspect that in our YIG|Pt microdisk, the splitting of modes is not sufficient to prevent nonlinear interactions that limit the amplitude of auto-oscillations [11]. In order to favor larger auto-oscillation amplitudes, YIG structures that are even more confined laterally (below 1 μ m) should be used [27], or one should excite a bullet mode [13].

In conclusion, we have demonstrated that it is possible to control electronically the SW damping in a YIG microdisk. Extending this result to one-dimensional SW guide [44] will offer great prospect in the emerging field of magnonics [45, 46], whose aim is to investigate the manipulation of SW and their quanta – magnons – with the benefice of combining ultra-low energy consumption and compactness. To improve the magnonic paradigm, a solution will be to actively compensate damping in the YIG magnetic insulator by SW amplification through stimulated emission generated by a charge current in the adjacent metallic layer with strong SOI.

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